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A new anomalous contribution to the central charge of the N=2 monopole

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ABSTRACT

We calculate the one-loop corrections to the mass and central charge of the BPS monopole in $N = 2$ super-Yang-Mills theory in 3+1 dimensions using a supersymmetry-preserving version of dimensional regularization adapted to solitons. In the renormalization scheme where previous studies have indicated vanishing quantum corrections, we find nontrivial corrections that we identify as the 3+1 dimensional analogue of the anomaly in the conformal central charge of the $N = 1$ supersymmetric kink in 1+1 dimensions. As in the latter case, the associated contribution to the ordinary central charge has exactly the required magnitude to preserve BPS saturation at the one-loop level. It also restores consistency of calculations involving sums over zero-point energies with the low-energy effective action of Seiberg and Witten.

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Supersymmetric (susy) solitons which saturate the Bogomolnyi bound [1, 2] have been found in various models and in various dimensions, and they play an important role in recent studies of nonperturbative effects in (susy) field theory and of duality [3].

In the earliest calculation of quantum corrections to these solitons it was assumed that supersymmetry would ensure complete cancellation of quantum corrections [4], thereby trivially guaranteeing Bogomolnyi-Prasad-Sommerfield (BPS) saturation [1, 5] at the quantum level. However, more careful calculations found that there may be nonzero corrections already in the simplest example of the 1+1-dimensional susy kink, but by the end of the 1980's the literature was in unresolved disagreement concerning the correct value. This issue was reopened in 1997 when two of us [6] noted that some of the earlier calculations had used methods which gave wrong results when applied to the exactly solvable bosonic sine-Gordon kink. The result for the mass obtained in [6] was however contaminated by energy located at the boundary of the quantization volume, which was corrected subsequently in [7] by the use of topological boundary conditions. This singled out as correct the earlier result of Schonfeld [8], who considered a kink-antikink system, as well as of Casahorrán [9], who used a finite mass formula in terms of only the discrete modes [10], and refuted the null results of Refs. [11, 12, 13, 14]. However it led to a new problem because it seemed to imply a violation of the Bogomolnyi bound, since the central charge did not appear to receive quantum corrections [15].

Ref. [7] suspected an anomaly at work, and shortly thereafter Shifman, Vainshtein, and Voloshin [16] showed that supersymmetry enforces an anomalous contribution to the central charge which is in fact part of a susy multiplet involving other better-known anomalies, the trace and conformal-susy anomaly.

In Ref. [17] we recently developed a version of dimensional regularization which preserves susy and reproduces the correct susy kink mass without the complications of other methods. In Ref. [18] we then demonstrated how the anomalous contribution to the central charge of the susy kink can be obtained as a remnant of parity violation in the odd-dimensional model used for embedding the susy kink. We also showed that the anomaly in question is not an anomaly in the ordinary central charge, but in the conformal central charge. The ordinary central charge itself has no anomaly, because it is produced by the anticommutator of two ordinary susy charges and the latter are free from anomalies. However, we found that one can define a confor-

mal central-charge current whose divergence is proportional to the ordinary central-charge current, and the anomaly in the former is directly related to nontrivial quantum corrections to the latter.

Most recently, we applied this method to the $N = 2$ vortex in 2+1 dimensions [19] and demonstrated BPS saturation despite nonvanishing quantum corrections. However, in this case these quantum corrections are in no way related to an anomaly (in the conformal central charge), as there is also no trace and conformal-susy anomaly in odd dimensions. Nevertheless, the corrections to the central charge turned out to be nontrivial in that one has to take into account the effect of the background vortex on the quantum fluctuations far away from the vortex. The value of the mass correction agreed with a result deduced by heat-kernel methods [20].

In this Letter, we consider the $N = 2$ monopole in 3+1 dimensions, which has been used by many authors in studies of duality. The monopole model has more unbroken susy generators than the susy kink or the vortex, so one runs the risk (or the blessing) of vanishing quantum corrections. This model has been studied before in Refs. [21, 22, 23] and while the initial result of vanishing corrections of Ref. [21] turned out to be an oversimplification, Refs. [22, 23] nevertheless arrived at the conclusion of vanishing quantum corrections, at least in the simplest renormalization scheme.

By setting up a suitable susy-preserving dimensional regularization method which embeds the monopole into 4+1 dimensions, we shall verify explicitly that BPS saturation holds[1], but we find a nonvanishing contribution to the mass (in the simplest renormalization scheme), which is matched by an anomalous contribution to the central charge operator coming from parity-violating quantum corrections to the additional momentum component, precisely analogous to the situation in the susy kink.

The $N = 2$ super-Yang-Mills theory in 3+1 dimensions can be obtained by dimensional reduction from the (5+1)-dimensional $N = 1$ theory [24]

$$\mathcal{L} = -\frac{1}{4} F_{AB}^2 - \bar{\lambda}\Gamma^A D_A \lambda, \tag{1}$$

where the indices A, B take the values $0, 1, 2, 3, 5, 6$ and which is invariant under

$$\delta A_B^a = \bar{\lambda}^a \Gamma_B \eta - \bar{\eta}\Gamma_B \lambda^a, \quad \delta\lambda^a = \frac{1}{2} F_{BC}^a \Gamma^B \Gamma^C \eta. \tag{2}$$

[1] Since BPS saturation is guaranteed by the multiplet shortening arguments of Ref. [2], this just verifies that our regularization method indeed preserves susy.

The complex spinor λ is in the adjoint representation of the gauge group which we assume to be SU(2) in the following and $(D_A\lambda)^a = (\partial_A\lambda + gA_A\times\lambda)^a = \partial_A\lambda^a + g\epsilon^{abc}A_A^b\lambda^c$. Furthermore, λ and η satisfy the Weyl condition:

$$(1-\Gamma_7)\lambda = 0 \quad \text{with} \quad \Gamma_7 = \Gamma_0\Gamma_1\Gamma_2\Gamma_3\Gamma_5\Gamma_6. \tag{3}$$

To carry out the dimensional reduction we write $A_B = (A_\mu, P, S)$ and choose the following representation of gamma matrices

$$\begin{aligned} \Gamma_\mu &= \gamma_\mu\otimes\sigma_1 \quad , \quad \mu = 0,1,2,3, \\ \Gamma_5 &= \gamma_5\otimes\sigma_1 \quad , \quad \Gamma_6 = 1\!\!1\otimes\sigma_2. \end{aligned} \tag{4}$$

In this representation the Weyl condition (3) becomes $\lambda = \binom{\psi}{0}$, with a complex four-component spinor ψ.[2]

The (3+1)-dimensional Lagrangian then reads

$$\begin{aligned} \mathcal{L} &= -\{\frac{1}{4}F_{\mu\nu}^2 + \frac{1}{2}(D_\mu S)^2 + \frac{1}{2}(D_\mu P)^2 + \frac{1}{2}g^2(S\times P)^2\} \\ &\quad -\{\bar\psi\gamma^\mu D_\mu\psi + ig\bar\psi(S\times\psi) + g\bar\psi\gamma_5(P\times\psi)\}. \end{aligned} \tag{5}$$

We choose the symmetry-breaking Higgs field as $S^a \equiv A_6^a = v\delta_3^a$ in the trivial sector. The BPS monopoles are of the form (for $A_0 = 0$) [5]

$$A_i^a = \epsilon_{aij}\frac{x^j}{gr^2}(1-K(mr)), \tag{6}$$

$$S^a = \delta_i^a\frac{x^i}{gr^2}H(mr), \tag{7}$$

with $H = mr\coth(mr) - 1$ and $K = mr/\sinh(mr)$, where $m = gv$ is the mass of the particles that are charged under the unbroken U(1). The BPS equation $F_{ij}^a + \epsilon_{ijk}D_kS^a = 0$ can be written as a self-duality equation for F_{MN} with $M,N = 1,2,3,6$, and the classical mass is $M_{\text{cl.}} = 4\pi m/g^2$.

The susy algebra for the charges $Q^\alpha = \int j^{0\alpha}d^3x$ with $j^A = \frac{1}{2}\Gamma^B\Gamma^C F_{BC}\Gamma^A\lambda$ reads

$$\{Q^\alpha, \bar Q_\beta\} = -(\gamma^\mu)^\alpha{}_\beta P_\mu + (\gamma_5)^\alpha{}_\beta\, U + i\delta_\beta^\alpha\, V, \tag{8}$$

[2]We use the metric with signature $(-,+,+,+,+,+)$ and $\bar\lambda^a = (\lambda^a)^\dagger i\Gamma^0$, hence $\bar\psi^a = (\psi^a)^\dagger i\gamma^0$. One can rewrite this model in terms of two symplectic Majorana spinors in order to exhibit the R symmetry group U(2).

with $\alpha, \beta = 1, \ldots, 4$. In the trivial sector P_μ acts as ∂_μ, and U multiplies the massive fields by m, but in the topological sector P_μ are covariant translations, and U and V are surface integrals. The classical monopole solution saturates the BPS bound $M^2 \geq |\langle U \rangle|^2 + |\langle V \rangle|^2$ by $|U_{\rm cl.}| = M_{\rm cl.}$, and $V_{\rm cl.} = 0$. Quantum corrections may change the values of matrix elements of the charges in (8), but the algebra (8) remains unmodified and the charges themselves conserved, i.e., without anomalies, as we discussed above.

For obtaining the one-loop quantum corrections, one has to consider quantum fluctuations about the monopole background. The bosonic fluctuation equations turn out to be simplest in the background-covariant Feynman-R_ξ gauge which is obtained by dimensional reduction of the ordinary background-covariant Feynman gauge-fixing term in (5+1) dimensions $-\frac{1}{2}(D_B[\hat{A}]\, a^B)^2$, where a^B comprises the bosonic fluctuations and $\hat{A}^B$ the background fields. As has been found in Refs. [22, 23], in this gauge the eigenvalues of the bosonic fluctuation equations (taking into account Faddeev-Popov fields) and those of the fermionic fluctuation equations combine such that one can make use of an index-theorem by Weinberg [25] to determine the spectral density. This leads to the following (unregularized!) formula for the one-loop mass correction

$$M^{(1)} = \frac{4\pi m_0}{g_0^2} + \frac{\hbar}{2} \sum (\omega_B - \omega_F) = \frac{4\pi m_0}{g_0^2} + \frac{\hbar}{2} \int \frac{d^3k}{(2\pi)^3} \sqrt{k^2 + m^2}\, \rho_M(k^2), \tag{9}$$

with m_0 and g_0 denoting bare quantities and

$$\rho_M(k^2) = \frac{-8\pi m}{k^2(k^2 + m^2)}. \tag{10}$$

This expression is logarithmically divergent and is made finite by combining it with the one-loop renormalization of g, while m does not need to be renormalized [22, 23].

Defining renormalized quantities in the trivial sector and using the background-covariant Feynman-R_ξ gauge one finds that tadpole diagrams cancel among themselves. Because of background-covariance it suffices to formulate a renormalization condition for one of the two-point functions, and a particularly simple and natural choice is to renormalize the two-point functions of the massless bosons on-shell. With such a choice, Refs. [22, 23] came to the conclusion that the counterterms precisely cancel the contribution from the

zero-point energies[3].

However, while [22] did not specify the regularization method used to obtain this result, Ref. [23] regularized by inserting slightly different oscillatory factors in the two-point function and in the integral over the spectral density ρ_M, a procedure that is not obviously self-consistent.

We shall instead use dimensional regularization in a supersymmetry preserving manner, namely by embedding the (3+1)-dimensional theory and the BPS monopole in up to one higher dimension (the x^5 direction), where the BPS monopole can be trivially extended into a string-like object. For the purpose of dimensional regularization of the (3+1)-dimensional model this is sufficient; it corresponds to trivial Kaluza-Klein reduction of x^6, and continuous dimensional reduction from (4+1) to (3+1) dimensions.

The so dimensionally regularized mode sum can then be written as

$$\frac{\hbar}{2}\sum(\omega_B-\omega_F)=\frac{\hbar}{2}\int\frac{d^3k\,d^\epsilon\ell}{(2\pi)^{3+\epsilon}}\sqrt{k^2+\ell^2+m^2}\,\rho_M(k^2) \tag{11}$$

with ρ_M still given by (10).

At this point we have a choice how to present our results: we can either show BPS saturation (and its nontrivial ingredients) at the unrenormalized (but regularized!) level and thus remain independent of specific renormalization prescriptions, or we can first renormalize the theory and give well-defined final results also. Since the 'minimal' renormalization scheme introduced above is the most widely used one and since therein previous works have obtained null results, we opt for the latter and just remark that BPS saturation itself as well as the anomalous contribution that we shall derive are both independent of the details of the renormalization procedure.

Renormalizing the on-shell photon self-energy in background covariant $R_{\xi=1}$ gauge one obtains

$$\frac{1}{g_0^2}=\frac{1}{g^2}+4\hbar\int\frac{d^{4+\epsilon}}{(2\pi)^{4+\epsilon}}\frac{1}{(k_E^2+m^2)^2} \tag{12}$$

where the index E in k_E refers to Euclidean signature.

[3]Ref. [23] also considered other renormalization schemes and other gauge choices, where there are finite remainders and left the question of existence of quantum corrections to the monopole mass to some extent open. In the present paper we shall restrict ourselves to discussing the above "minimal" scheme.

Now, carrying out the ℓ-integration in the sum over zero-point energies (11) gives (setting from now on $\hbar = 1$)

$$\frac{1}{2}\sum(\omega_B - \omega_F) = -\frac{2m}{\pi}\frac{\Gamma(-\frac{1}{2}-\frac{\epsilon}{2})}{(2\pi^{\frac{1}{2}})^{\epsilon}\Gamma(-\frac{1}{2})}\int_0^\infty dk(k^2+m^2)^{-\frac{1}{2}+\frac{\epsilon}{2}}, \tag{13}$$

while $3+\epsilon$ integrations in the counterterm in (12) yield

$$\delta M \equiv 4\pi m(g_0^{-2} - g^{-2}) = 16\pi m\frac{\Gamma(\frac{1}{2}-\frac{\epsilon}{2})}{(2\pi^{\frac{1}{2}})^{3+\epsilon}}\int_{-\infty}^{\infty}\frac{dk_4}{2\pi}(k_4^2+m^2)^{-\frac{1}{2}+\frac{\epsilon}{2}}. \tag{14}$$

Combining these two expressions we find that there is a mismatch proportional to ϵ, but ϵ multiplies a logarithmically divergent integral, which in dimensional regularization involves a pole ϵ^{-1}. We therefore obtain a finite correction of the form

$$\begin{aligned} M^{(1)} &= \frac{4\pi m}{g^2} - \epsilon\times\frac{2m}{\pi}\frac{\Gamma(-\frac{1}{2}-\frac{\epsilon}{2})}{(2\pi^{\frac{1}{2}})^{\epsilon}\Gamma(-\frac{1}{2})}\int_0^\infty dk(k^2+m^2)^{-\frac{1}{2}+\frac{\epsilon}{2}} \\ &= \frac{4\pi m}{g^2} - \frac{2m}{\pi} + O(\epsilon) \end{aligned} \tag{15}$$

which because of the fact that it arises as $0\times\infty$ bears the fingerprint of an anomaly.

Indeed, as we shall now show, this result is completely analogous to the case of the $N=1$ susy kink in (1+1) dimensions, where a nonvanishing quantum correction to the kink mass (in a minimal renormalization scheme) is associated with an anomalous contribution to the central charge (which is scheme-independent; in a non-minimal renormalization scheme there are also non-anomalous quantum corrections to the central charge).

In Ref. [23] it has been argued that in the renormalization scheme defined above, the one-loop contributions to the central charge precisely cancel the contribution from the counterterm in the classical expression. In this particular calculation it turns out that the cancelling contributions have identical form so that the regularization methods of Ref. [23] can be used at least self-consistently, and also straightforward dimensional regularization would imply complete cancellations. The result (15) would then appear to violate the Bogomolnyi bound.

However, this is just the situation encountered in the (1+1)-dimensional susy kink. As we have shown in Ref. [18], dimensional regularization gives

a zero result for the correction to the central charge unless the latter is augmented by the momentum operator in the extra dimension used to embed the soliton. This is necessary for manifest supersymmetry, and, indeed, the extra momentum operator can acquire a nonvanishing expectation value. As it turns out, the latter is entirely due to nontrivial contributions from the fermions $\psi = \binom{\psi_+}{\psi_-}$, whose fluctuation equations have the form

$$
\begin{aligned}
L\psi_+ + i(\partial_t + \partial_5)\psi_- &= 0, && (16)\\
i(\partial_t - \partial_5)\psi_+ + L^\dagger\psi_- &= 0. && (17)
\end{aligned}
$$

The fermionic field operator can be written as

$$
\begin{aligned}
\psi(x) &= \int \frac{d^\epsilon \ell}{(2\pi)^{\epsilon/2}} \sum\!\!\!\!\!\!\int \frac{d^3k}{(2\pi)^{3/2}} \frac{1}{\sqrt{2\omega}} \Big\{ a_{kl} e^{-i(\omega t - \ell x^5)} \begin{pmatrix} \sqrt{\omega - \ell}\, \chi_+ \\ -\sqrt{\omega + \ell}\, \chi_- \end{pmatrix} \\
&\qquad + b^\dagger_{kl} e^{i(\omega t - \ell x^5)} \begin{pmatrix} \sqrt{\omega - \ell}\, \chi_+ \\ \sqrt{\omega + \ell}\, \chi_- \end{pmatrix} \Big\} \qquad (18)
\end{aligned}
$$

where $\chi_- = \frac{1}{\omega_k} L\chi_+$ and $\chi_+ = \frac{1}{\omega_k} L^\dagger \chi_-$ with $\omega_k^2 = k^2 + m^2$, and the normalization factors $\sqrt{\omega \pm \ell}$ are such that $L^\dagger L\chi_+ = \omega^2\chi_+$ and $LL^\dagger\chi_- = \omega^2\chi_-$ with $\omega^2 = \omega_k^2 + \ell^2$. Because of these normalization factors, one obtains an expression for the momentum density Θ_{05} in the extra dimension which has an even component under reflection in the extra momentum variable ℓ

$$
\begin{aligned}
\langle \Theta_{05} \rangle &= \int \frac{d^\epsilon \ell}{(2\pi)^\epsilon} \int \frac{d^3k}{(2\pi)^3} \frac{\ell}{2\omega} \left[(\omega - \ell)|\chi_+|^2 + (\omega + \ell)|\chi_-|^2 \right] \\
&= \int \frac{d^\epsilon \ell}{(2\pi)^\epsilon} \int \frac{d^3k}{(2\pi)^3} \frac{\ell^2}{2\omega} (|\chi_-|^2 - |\chi_+|^2) \qquad (19)
\end{aligned}
$$

(omitting zero-mode contributions which do not contribute in dimensional regularization [18]).

Integration over x then produces the spectral density (10) and finally yields

$$
\begin{aligned}
\Delta U_{\rm an} &= \int d^3x\, \langle \Theta_{05} \rangle = \int \frac{d^3k\, d^\epsilon \ell}{(2\pi)^{3+\epsilon}} \frac{\ell^2}{2\sqrt{k^2 + \ell^2 + m^2}} \rho_M(k^2) \\
&= -4m \int_0^\infty \frac{dk}{2\pi} \int \frac{d^\epsilon \ell}{(2\pi)^\epsilon} \frac{\ell^2}{(k^2 + m^2)\sqrt{k^2 + \ell^2 + m^2}} \\
&= -8 \frac{\Gamma(1 - \frac{\epsilon}{2})}{(4\pi)^{1+\frac{\epsilon}{2}}} \frac{m^{1+\epsilon}}{1+\epsilon} = -\frac{2m}{\pi} + O(\epsilon), \qquad (20)
\end{aligned}
$$

which is indeed equal to the nonzero mass correction obtained above.

We thus have verified that the BPS bound remains saturated under quantum corrections, but the quantum corrections to mass and central charge both contain an anomalous contribution, analogous to the anomalous contribution to the central charge of the 1+1 dimensional minimally supersymmetric kink.

The nontrivial result (20) is in fact in complete accordance with the low-energy effective action for $N = 2$ super-Yang-Mills theory as obtained by Seiberg and Witten [3].[4] According to the latter, the low-energy effective action is fully determined by a prepotential $\mathcal{F}(A)$, which to one-loop order is given by

$$\mathcal{F}_{1-\mathrm{loop}}(A) = \frac{i}{2\pi} A^2 \ln \frac{A^2}{\Lambda^2}, \tag{21}$$

where A is a chiral superfield and Λ the scale parameter of the theory generated by dimensional transmutation. The value of its scalar component a corresponds in our notation to $gv = m$. In the absence of a θ parameter, the one-loop renormalized coupling is given by

$$\frac{4\pi i}{g^2} = \tau(a) = \frac{\partial^2 \mathcal{F}}{\partial a^2} = \frac{i}{\pi} \left(\ln \frac{a^2}{\Lambda^2} + 3 \right). \tag{22}$$

This definition agrees with the "minimal" renormalization scheme that we have considered above, because the latter involves only the zero-momentum limit of the two-point function of the massless fields. For a single magnetic monopole, the central charge is given by

$$|U| = a_D = \frac{\partial \mathcal{F}}{\partial a} = \frac{i}{\pi} a \left(\ln \frac{a^2}{\Lambda^2} + 1 \right) = \frac{4\pi a}{g^2} - \frac{2a}{\pi}, \tag{23}$$

and since $a = m$, this exactly agrees with the result of our direct calculation in (20).

Now, the low-energy effective action associated with (21) has been derived from a consistency requirement with the anomaly of the $\mathrm{U}(1)_R$ symmetry of the microscopic theory. The anomaly in the conformal central charge, which we have identified as being responsible for the entire nonzero correction (20), is evidently consistent with the former. Just as in the case of the minimally supersymmetric kink in 1+1 dimensions, it constitutes a new anomaly[5] that

[4] We are grateful to Horatiu Nastase for pointing this out to us.

[5] The possibility of anomalous contributions to the central charges of $N = 2$ super-Yang-Mills theories in 4 dimensions has most recently also been noted in [26], however without a calculation of the coefficients.

had previously been missed in direct calculations [22, 23] of the quantum corrections to the $N = 2$ monopole using sums over zero-point energies.

We intend to discuss the details of the anomalous conformal central-charge current further in a future publication. We also plan then to consider dimensional regularization by dimensional reduction starting from 3+1 dimensions (instead of from 4+1 dimensions as we have done in this paper), which, as we have shown in Ref. [18], locates the anomaly in an evanescent counterterm.

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